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Effect of Spin Current on Uniform Ferromagnetism: Domain Nucleation

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Large spin current applied to a uniform ferromagnet leads to a spin-wave instability as pointed out recently. In this paper, it is shown that such spin-wave instability is absent in a state containing a domain wall, which indicates that nucleation of magnetic domains occurs above a certain critical spin current. This scenario is supported also by an explicit energy comparison of the two states under spin current.

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Magnetization dynamics driven by spin torque from a spin polarized current (spin current) has been studied extensively after the theoretical prediction [1, 2] that a spin current can be used to flip the magnetization in pillar (or spin valve) structures [3, 4, 5, 6, 7, 8]. As a theoretical framework to describe such current-induced magnetization dynamics, Bazaliy, Jones and Zhang (BJZ) [9] derived a modified Landau-Lifshitz-Gilbert (LLG) equation for a fully-polarized ferromagnet (half metal) with a new term proportional to $\boldsymbol{j}\cdot\nabla\boldsymbol{n}$. This new term represents a spin torque that a current of density $\boldsymbol{j}$ exerts on a local magnetization $\boldsymbol{n}$. Later on, it was generalized to the partially polarized case with an interpretation of $\boldsymbol{j}$ as the spin current density, $\boldsymbol{j}_{\rm s}$ [10, 11]. In the Hamiltonian, the effect of spin torque appears in a form

$$H_{\rm ST} = \int d^3x \frac{\hbar}{2e}\boldsymbol{j}_{\rm s}\cdot\nabla\phi(1-\cos\theta), \tag{1}$$

where θ and ϕ are polar angles which parameterize $\boldsymbol{n}$. As seen from this form, spin current favors a magnetic configuration with spatial gradient, or more precisely, with finite Berry-phase curvature. It is thus expected that a large spin current destabilizes a uniform ferromagnetic state. This is indeed seen from the spin-wave energy [9, 10, 11],

$$\Omega^{\rm uni}_{\boldsymbol{k}} = \frac{KS}{\hbar}\sqrt{(k^2\lambda^2+1)(k^2\lambda^2+1+\kappa)} + \boldsymbol{k}\cdot\boldsymbol{v}_{\rm s}, \tag{2}$$

where $\lambda = \sqrt{J/K}$ and $\kappa = K_\perp/K$, with J, K and $K_\perp$ being exchange constant, easy- and hard-axis anisotropy constants, respectively, of localized spins. The first term is the well-known spin-wave dispersion with anisotropy gap. The effect of spin current appears in the second term as the "Doppler shift" [10], where $\boldsymbol{v}_{\rm s}(\propto \boldsymbol{j}_{\rm s})$ represents drift velocity of electron spins. For sufficiently large $\boldsymbol{v}_{\rm s}$, $\Omega^{\rm uni}_{\boldsymbol{k}}$ becomes negative for a range of $\boldsymbol{k}$. This means that there exist states with negative excitation energy, indicating the instability of the assumed uniformly-magnetized state [9, 10, 11]. The true ground state under spin current, however, remains to be identified.

In this Letter, we point out that possible ground state is a state containing domain walls. Namely, the energy of a domain wall becomes lower than the uniform ferromagnetic state when a spin current exceeds a certain critical value $j_{\rm s}^{\rm cr}$. The behavior of nucleated domains depends much on the magnetic anisotropy parameters; Flowing domain wall is stable when the anisotropy is of uniaxial type. For the case of large hard-axis anisotropy, static domain wall is stable at the nucleation threshold, but when it starts to flow under larger current, even a state with a domain wall becomes unstable. The nature of the resulting state is still unknown.

Our prediction of domain formation by spin current may be related to the very recent experimental observations in metallic and semiconducting pillars and films [12, 13, 14, 15, 16, 17], which suggest spatially inhomogeneous magnetization reversal. Domain nucleation is also suggested by recent numerical simulation [18].

We start by extending the formulation of BJZ [9] to arbitrary degree of polarization, and derive the effective Lagrangian for magnetization which is slowly varying in space and time by assumption. With this effective Lagrangian, we calculate spin-wave dispersion around a domain wall solution in the presence of spin current. The spin-wave Doppler shift term now has the form, $k(v_{\rm s}-\dot{X})$, where $\dot{X}$ is the speed of the domain wall. From this re-

sult, we find that the spin-wave instability does not occur around a domain wall for any large spin current if the anisotropy is of uniaxial type ($K > K_\perp$). This suggests that *the true ground state is a multi-domain state with domain walls*. These domain walls turn out to be flowing with average speed equal to drift velocity v_s of electron spin determined by the spin current. The stability is then understood from Galilean invariance, since a domain wall moving with speed $\dot{X} = v_\mathrm{s}$ is equivalent to a domain wall at rest and in the absence of spin current.

The situation is slightly different in the case of large hard-axis anisotropy, $K_\perp >> K$. Domain walls created by the spin current above j_s^cr are at rest if the current is below depinning threshold, $j_\mathrm{s}^\mathrm{depin} \propto K_\perp\lambda (> j_\mathrm{s}^\mathrm{cr})$. At higher current, the walls start to flow, but this triggers another spin-wave instability. Thus multi-domain state collapses for $j_\mathrm{s} > j_\mathrm{s}^\mathrm{depin}$ into another new ground state which is still unknown [19]. We defer this new state to future studies.

We consider a ferromagnet consisting of localized spins and conduction electrons. The spins are assumed to have an easy z-axis and a hard y-axis, and described by the Lagrangian

$$L_\mathrm{S} = \int \frac{d^3x}{a^3} \hbar S \dot{\phi}(\cos\theta - 1) - H_\mathrm{S}, \quad (3)$$

$$H_\mathrm{S} = \int \frac{d^3x}{a^3} \left\{ \frac{JS^2}{2}(\nabla \boldsymbol{n})^2 + \frac{S^2}{2}\sin^2\theta\left(K + K_\perp \sin^2\phi\right) \right\}. \quad (4)$$

We have adopted a continuum description for localized spins, $\boldsymbol{S}(\boldsymbol{x},t) = S\,\boldsymbol{n}(\boldsymbol{x},t)$, with unit vector $\boldsymbol{n}(\boldsymbol{x},t) = (\sin\theta\cos\phi, \sin\theta\sin\phi, \cos\theta)$ and the magnitude of spin, S. The J is the exchange constant, and a the lattice constant. The easy-axis (K) and hard-axis ($K_\perp$) anisotropy constants generally incorporate the effect of demagnetizing field. The exchange interaction between localized spins and conduction electrons is given by

$$\hat{H}_\mathrm{sd} = -\Delta \int d^3x\, \boldsymbol{n}(\boldsymbol{x},t)\cdot(\hat{c}^\dagger(\boldsymbol{x},t)\boldsymbol{\sigma}\hat{c}(\boldsymbol{x},t)), \quad (5)$$

where 2Δ and $\hat{c}$ ($\hat{c}^\dagger$) are the energy splitting and annihilation (creation) operator of conduction electrons, respectively, and $\boldsymbol{\sigma}$ is a Pauli-matrix vector. The free-electron part is given by $\hat{H}_\mathrm{el}^0 = \sum_{\boldsymbol{k}\sigma} \epsilon_{\boldsymbol{k}} \hat{c}_{\boldsymbol{k}\sigma}^\dagger \hat{c}_{\boldsymbol{k}\sigma}$ with $\epsilon_{\boldsymbol{k}} = \hbar^2\boldsymbol{k}^2/2m$. We perform a local gauge transformation [20, 21] in electron spin space so that the quantization axis is parallel to $\boldsymbol{S}(\boldsymbol{x},t)$ at each point of space and time; $\hat{c}(\boldsymbol{x},t) = U(\boldsymbol{x},t)\hat{a}(\boldsymbol{x},t)$, where $\hat{a}(\boldsymbol{x},t)$ is the two-component electron operator in the rotated frame, and $U(\boldsymbol{x},t) \equiv \boldsymbol{m}(\boldsymbol{x},t)\cdot\boldsymbol{\sigma}$ is an SU(2) matrix with $\boldsymbol{m}(\boldsymbol{x},t) = (\sin(\theta/2)\cos\phi, \sin(\theta/2)\sin\phi, \cos(\theta/2))$. The electron part of the Hamiltonian is now given by $\hat{H}_\mathrm{el} = \sum_{\boldsymbol{k}\sigma} \epsilon_{\boldsymbol{k}\sigma} a_{\boldsymbol{k}\sigma}^\dagger a_{\boldsymbol{k}\sigma}$ with $\epsilon_{\boldsymbol{k}\sigma} = \hbar^2\boldsymbol{k}^2/2m - \sigma\Delta$, and the interaction Hamiltonian $\hat{H}_\mathrm{int}$ by

$$\hat{H}_\mathrm{int} = \frac{\hbar}{V}\sum_{\boldsymbol{kq}\alpha\beta}\left\{A_0^\nu(\boldsymbol{q},t) + \frac{\hbar(2k_j+q_j)}{2m}A_j^\nu(\boldsymbol{q},t)\right\} \times \hat{a}_{\boldsymbol{k}+\boldsymbol{q}\alpha}^\dagger \sigma_{\alpha\beta}^\nu \hat{a}_{\boldsymbol{k}\beta} + \frac{\hbar^2}{2mV^2}\sum_{\boldsymbol{kqp}\alpha} A_j^\mu(\boldsymbol{p},t)A_j^\mu(\boldsymbol{q}-\boldsymbol{p},t)\hat{a}_{\boldsymbol{k}+\boldsymbol{q}\alpha}^\dagger \hat{a}_{\boldsymbol{k}\alpha} \quad (6)$$

Here, $\boldsymbol{A}_\mu(\boldsymbol{q},t) = \int d^3x e^{-i\boldsymbol{q}\cdot\boldsymbol{x}}\boldsymbol{A}_\mu(\boldsymbol{x},t)$ with $\boldsymbol{A}_\mu(\boldsymbol{x},t) = \boldsymbol{m}\times\partial_\mu\boldsymbol{m}$ represents SU(2) gauge field with space ($\mu = x,y,z$) and time ($\mu = 0$) components.

For slowly varying magnetic configurations, the effective Lagrangian can be derived by a perturbative expansion with respect to $\boldsymbol{A}_\mu$. This is the gradient expansion. In conformity with L_S, we retain terms up to second order in spatial gradient or in $\boldsymbol{A}_i$ ($i = x,y,z$), and first order in time derivative or in $\boldsymbol{A}_0$. The spin-torque term, $\sum_{\boldsymbol{x}} \frac{\hbar}{e}\boldsymbol{j}_\mathrm{s}\cdot\boldsymbol{A}^z$, arises from the first-order contribution $\langle\hat{H}_\mathrm{int}\rangle$. Here $\boldsymbol{j}_\mathrm{s} = \frac{e}{V}\sum_{\boldsymbol{k}\sigma}\sigma\frac{\hbar\boldsymbol{k}}{m}f_{\boldsymbol{k}\sigma}$ is the spin current density in the rotated frame, with a distribution function $f_{\boldsymbol{k}\sigma} = \langle\hat{a}_{\boldsymbol{k}\sigma}^\dagger(t)\hat{a}_{\boldsymbol{k}\sigma}(t)\rangle$ specifying the state carrying current. Other terms just give renormalizations of J and S [20]. The effective Lagrangian is thus given by

$$L_\mathrm{eff} = L_\mathrm{S} - H_\mathrm{ST}. \quad (7)$$

The second term, given by Eq. (1), with j_s is identical to that derived by BJZ [9] for the half-metallic case. It is seen that spin current favors a finite Berry phase curvature along the current. The modified LLG equation can be obtained by taking variations of L_eff with respect to θ and ϕ. It has the same form as in BJZ, and is applicable to general slowly-varying spin texture, and to the case of arbitrary degree of polarization.

Let us now study the spin-wave excitation around a domain-wall solution under spin current by using the method of collective coordinate [22]. The wall is assumed to be planar and moving in the direction perpendicular to the plane (chosen as x-direction). It is convenient to use a complex field $\xi(x,t) = e^{i\phi(x,t)}\tan\frac{\theta(x,t)}{2}$ to represent magnetization. We decompose it into the domain-wall configuration and spin waves around it:

$$\xi(x,t) = \xi_\mathrm{dw}(x,t)\exp\left[2\cosh\{(x-X)/\lambda\} \times \frac{1}{\sqrt{L}}\sum_k \varphi_k(x-X)\eta_k(t)\right], \quad (8)$$

where $\xi_\mathrm{dw} = \exp[-(x-X(t))/\lambda + i\phi_0(t)]$ represents a domain-wall configuration. The wall position $X(t)$ and the angle $\phi_0(t)$ between wall spins and the easy plane are proper collective coordinates [21, 23, 24, 25, 26]. The spin-wave part is expanded with modes $\varphi_k(x) = \frac{1}{\sqrt{k^2\lambda^2+1}}(-ik\lambda + \tanh\frac{x}{\lambda})e^{ikx}$ [23, 26], whose amplitude η_k precisely corresponds to the Holstein-Primakoff boson. (L is the length of the system in the x-direction.)

Substituting (8) into (7), we obtain the spin-wave Lagrangian $L_{\rm sw}$, up to the quadratic order in η, as

$$L_{\rm sw} = \frac{\hbar NS}{2\lambda}\sum_k \Big\{ i(\eta_k^*\dot{\eta}_k - \dot{\eta}_k^*\eta) \\ -2\left\{\Omega_k^{(0)} - k(\dot{X}-v_{\rm s})\right\}\eta_k^*\eta_k \\ +\frac{v_{\rm c}}{\lambda}(1-4\sin^2\phi_0)(\eta_{-k}\eta_k+\eta_{-k}^*\eta_k^*)\Big\}, \quad (9)$$

where $N = 2\lambda A/a^3$ is the number of spins in the wall (A is the cross-sectional area of the wall), $\Omega_k^{(0)} = \frac{KS}{\hbar}(k^2\lambda^2 + 1 + \kappa/2)$ and $v_{\rm c} = \frac{\lambda K_\perp S}{2\hbar}$. Details of the calculation will be presented in a separate paper. From Eq. (9), we find the spin-wave dispersion around a current-driven domain wall:

$$\Omega_k^{\rm dw} = \frac{KS}{\hbar}\Big\{(k^2\lambda^2+1)(k^2\lambda^2+1+\kappa) \\ + 2\kappa^2\sin^2\phi_0\cos 2\phi_0\Big\}^{1/2} + k\left(v_{\rm s}-\dot{X}\right). \quad (10)$$

The most important difference from the case of uniform ferromagnet is that the drift velocity $v_{\rm s} = \frac{a^3}{2eS}j_{\rm s}$ of spin current appears as a relative velocity, $v_{\rm s}-\dot{X}$, with respect to the moving wall.

Equation (10) shows that the spin-wave energy depends strongly on the wall dynamics (through $\dot{X}$ and ϕ_0). From the equations of motion for the domain wall [21], $\lambda\dot{\phi_0} + \alpha\dot{X} = 0$ and $\dot{X} - \alpha\lambda\dot{\phi_0} = v_{\rm c}\sin 2\phi_0 + v_{\rm s}$, where $\alpha \ll 1$ is a damping parameter, we have $(1+\alpha^2)\dot{X} = v_{\rm c}\sin 2\phi_0 + v_{\rm s}$. Neglecting terms of $\sim O(\alpha^2)$, we have,

$$\Omega_k^{\rm dw} = \frac{KS}{\hbar}\Big[\Big\{(k^2\lambda^2+1)(k^2\lambda^2+1+\kappa) \\ + 2\kappa^2\sin^2\phi_0\cos 2\phi_0\Big\}^{1/2} - \frac{k\lambda}{2}\kappa\sin 2\phi_0\Big], \quad (11)$$

which depends on $j_{\rm s}$ implicitly through ϕ_0.

For $K > K_\perp$, it is easy to see that $\Omega_k^{\rm dw}$ remains positive for all values of k and $j_{\rm s}$ [27]. Thus a domain-wall state is stable irrespective of the magnitude of the spin current in this case. This indicates that a uniform ferromagnetism collapses into a multi-domain configuration for $j_{\rm s} > j_{\rm s}^{\rm cr}$, where $j_{\rm s}^{\rm cr} = \frac{2eS^2}{\hbar a^3}K\lambda(1+\sqrt{1+\kappa})$ is the critical spin current density for the instability of a uniform ferromagnet [9, 10, 11]. Since the domain wall flows for $j_{\rm s} > j_{\rm s}^{\rm depin} \equiv \frac{eS^2}{\hbar a^3}K\lambda\kappa$ [21], which satisfies $j_s^{\rm cr} > j_{\rm s}^{\rm depin}$, it starts to flow as soon as it is created (Fig.1(a)). The stability of moving domain wall state is natural from Galilean invariance since domain wall with velocity of $v_{\rm s}$ is equivalent to the static domain wall without spin current.

In the opposite case, $K_\perp > K$, the spin wave shows no anomaly as long as $j_{\rm s} < j_{\rm s}^{\rm depin}$, where the wall remains static with constant ϕ_0 $(= -\frac{1}{2}\sin^{-1}(v_{\rm s}/v_{\rm c}))$. Hence, for $j_{\rm s}^{\rm cr} < j_{\rm s} < j_{\rm s}^{\rm depin}$ ($j_{\rm s}^{\rm cr} < j_{\rm s}^{\rm depin}$ is realized when $K_\perp > 8K$), the static multi-domain state is realized. In contrast, as soon as the wall starts to move under larger spin current ($j_{\rm s} > j_{\rm s}^{\rm depin}$), the spin wave becomes unstable (Eq.(11)). Thus another ground state would appear. The wavelength of the unstable mode around the uniform magnetization, $k = -(K(K+K_\perp)/J^2)^{1/4}$[11], suggests that this unknown ground state has a spatial structure with short length scale of $\sim (KK_\perp/J^2)^{-1/4} = (K/K_\perp)^{1/4}\lambda < \lambda$.

We have thus shown that the spin-wave instability under spin current is avoided by the existence of domain wall. This may indicate that domain nucleation occurs under spin current. We cannot, however, answer here how many domains are created. The above argument holds for each segment larger than the wall thickness, λ, and thus the domain wall can be nucleated with spacing of $\sim \lambda$. For a correct estimate of the spacing, however, we need to consider the dipolar interaction among domains. We will not pursue this point in this letter.

We here present a supporting argument for the above scenario of domain formation by evaluating explicitly the total energy of a domain wall in the presence of spin current. We consider only the case $K \ll K_\perp$ and $j_{\rm s}^{\rm cr} < j_{\rm s} \ll j_{\rm s}^{\rm depin}$, where the wall can remain static and thus the energy comparison has physical meaning. A static domain wall configuration is given by $\theta(x,t) = \theta_{\rm s}(x) = 2\tan^{-1}\exp(-x/\lambda')$, $\phi(x,t) = \phi_{\rm s}(x)$, where $\lambda' = \lambda/\sqrt{1+\kappa\sin^2\phi_0}$ includes the effect of contraction of domain wall [28] ($\phi_0 \equiv -\frac{1}{2}\int dx\phi_s(\partial_x\theta_{\rm s})\sin\theta_{\rm s}$). Note that $\phi_{\rm s}(x)$ is not spatially constant when there is a spin current. For large hard-axis anisotropy, $K_\perp \gg \frac{\hbar a^3}{eS^2\lambda}j_{\rm s}$, substitution of the static solution $\theta_{\rm s}$ and $\phi_{\rm s}$ into the modified LLG equation leads to $\phi_0 \simeq -(\hbar a^3/2eS^2\lambda K_\perp)j_{\rm s} \ll 1$. The total energy, $H_{\rm S}+H_{\rm ST}$, of this static domain wall is then given by [29]

$$E_{\rm dw} = H_{\rm S}|_{\theta_{\rm s},\phi_{\rm s}} - \frac{\hbar}{2e}j_{\rm s}\int d^3x\phi_{\rm s}\partial_x\theta_{\rm s}\sin\theta_{\rm s} \\ \simeq NKS^2\left[1-\frac{1}{2}\left(\frac{\hbar a^3}{2eS^2\sqrt{KK_\perp}\lambda}j_{\rm s}\right)^2\right]. \quad (12)$$

Since the energy $E_{\rm uni}$ of the uniform state is zero independently of $j_{\rm s}$, we find that $E_{\rm dw} < E_{\rm uni}$ when $j_{\rm s} > \frac{2eS^2}{\hbar a^3}\sqrt{2KK_\perp}\lambda$ (Fig.1(b)). This critical value is of the same order of magnitude as $j_{\rm s}^{\rm cr}$ for $K_\perp \gg K$. Thus, the uniformly magnetized state is a 'false vacuum', which collapses into a multi-domain state via first order phase transition under spin current above a critical value.

The domain formation may be understood from the spin wave around uniform ferromagentism. In fact, the spin wave energy, $\Omega_{\boldsymbol{k}}^{\rm uni}$, given by Eq. (2), indicates that the typical wavelength of the spin wave excitation (determined by the minimum of the spin wave energy) is given by $k^{-1} \sim -\lambda^2(\sqrt{1+\kappa}KS)/(\hbar v_{\rm s})$ for small $j_{\rm s}$. As the current increases, the wavelength becomes shorter and

the critical current is given by the condition it becomes comparable to the wall thickness.

The above estimate of the critical current density obtained for a planar wall would be applicable also to the case of nucleation in films such as in Ref.[15], at least for an estimate of order of magnitude.

Let us estimate the critical current numerically. The magnitude of the current needed for nucleation in a wire is estimated as $j_{\rm s} \sim 3\times 10^{13}$ A/m^2 by using the materials parameters of Co: $a = 2.3$ Å, $S \sim 1$, $J = 3 \times 10^{-40}$ Jm^2, $K = 6.2 \times 10^{-24}$ J, ($\lambda = \sqrt{J/K} = 76$ Å), under the assumption that $K_\perp$ is mostly due to the demagnetizing field, $K_\perp = M_{\rm s}^2 a^3/\mu_0 = 3 \times 10^{-23}$ J. Recent experiment on a Co layer with a current through a point contact with domain suggests a domain wall formation for a maximum current density close to the contact region of $j \sim 5\times 10^{12-13}$ A/m^2 [13, 14, 15]. This value is in rough agreement with our estimate. Other recent experiments on metallic and semiconducting pillar structures indicating inhomogeneous magnetization reversal may also be due to domain (or domain-like) structure created by spin current [12, 16, 17].

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[1] J. C. Slonczewski, J. Magn. Magn. Mater. **159**, L1 (1996).
[2] L. Berger, Phys. Rev. B **54**, 9353 (1996).
[3] J. Z. Sun, J. Magn. Magn. Mater. **202**, 157 (1999).
[4] E. B. Myers, D. C. Ralph, J. A. Katine, R. N. Louie, and Buhrman, Science **285**, 867 (1999).
[5] J. A. Katine, F. J. Albert, R. A. Buhrman, E. B. Myers, and D. C. Ralph, Phys. Rev. Lett. **84**, 3149 (2000).
[6] J. Grollier, V. Cros, A. Hamzic, J. M. George, H. Jaffrès, A. Fert, G. Faini, J. Ben Youssef, and H. Legall, Appl. Phys. Lett. **78**, 3663 (2001).
[7] M. Tsoi, V. Tsoi, J. Bass, A. G. M. Jansen, and P. Wyder, Phys. Rev. Lett. **89**, 246803 (2002).
[8] S. I. Kiselev, J. C. Sankey, I. N. Krivorotov, N. C. Emley, R. J. Schoelkopf, R. A. Buhrman, and D. C. Ralph, Nature (London) **425**, 380 (2003).
[9] Ya. B. Bazaliy, B. A. Jones, and Shou-Cheng Zhang, Phys. Rev. B **57**, R3213 (1998).
[10] J. Fernández-Rossier, M. Braun, A.S. Núñez and A. H. MacDonald, Phys. Rev. B **69**, 174412 (2004).
[11] Z. Li and S. Zhang, Phys. Rev. Lett. **92**, 207203 (2004).
[12] J.-E. Wegrowe, D. Kelly, Y. Jaccard, Ph. Guittienne, and J.-Ph. Ansermet, Europhys. Lett. **45**, 626 (1999).
[13] Y. Ji, C. L. Chien, and M. D. Stiles, Phys. Rev. Lett. **90**, 106601 (2003).
[14] T. Y. Chen, Y. Ji, and C. L. Chien, Appl. Phys. Lett. **84**, 380 (2004).
[15] T. Y. Chen, Y. Ji, C. L. Chien, and M. D. Stiles, Phys. Rev. Lett. **93**, 026601 (2004).
[16] D. Chiba, Y. Sato, T. Kita, F. Matsukura, and H. Ohno, cond-mat/0403500.
[17] T. Yang, T. Kimura, and Y. Otani, preprint.
[18] A. Thiaville and Y. Nakatani, private communication.
[19] Numerical simulation observed dynamical and chaotic domain pattern [18].
[20] G. Tatara and H. Fukuyama Phys. Rev. Lett. **72**, 772 (1994); J. Phys. Soc. Jpn. **63**, 2538 (1994).
[21] G. Tatara and H. Kohno, Phys. Rev. Lett. **92**, 086601 (2004).
[22] R. Rajaraman, *Solitons and Instantons* (North-Holand, Amsterdam, 1982).
[23] H.-B. Braun and D. Loss, Phys. Rev. B**53**, 3237 (1996).
[24] S. Takagi and G. Tatara, Phys. Rev. B**54**, 9920 (1996).
[25] J. Shibata and S. Takagi, Phys. Rev. B**62**, 5719 (2000).
[26] J. Shibata and S. Takagi, Phys. Atom. Nucl. **64**, 2206 (2001).
[27] In fact, $\Omega_k^{\rm dw}$ vanishes when $f(x) \equiv (x+1)(x+1+\kappa) - \frac{\kappa^2}{4}x\sin^2 2\phi_0 + 2\kappa^2 \sin^2\phi_0\cos 2\phi_0 = 0$, where $x \equiv (k\lambda)^2 (\geq 0)$. At $x=0$, $f(0) = 1+\kappa+2\kappa^2\sin^2\phi_0\cos 2\phi_0$ and is positive if $\kappa < 1$ irrespective of ϕ_0. Noting that $f(x)$ in $x \geq 0$ is an increasing function of x for $(0 \leq)\kappa < 2(1+\sqrt{3})$, we see that $\Omega_k^{\rm dw} > 0$ if $\kappa < 1$.
[28] H.-B. Braun, J. Kyriakidis and D. Loss, Phys. Rev. B**56**, 8129 (1997).
[29] In this calculation, we have performed the integration by parts with respect to x in the spin-torque term of Eq.(7) before substituting the static solution, $\theta_{\rm s}$ and $\phi_{\rm s}$.

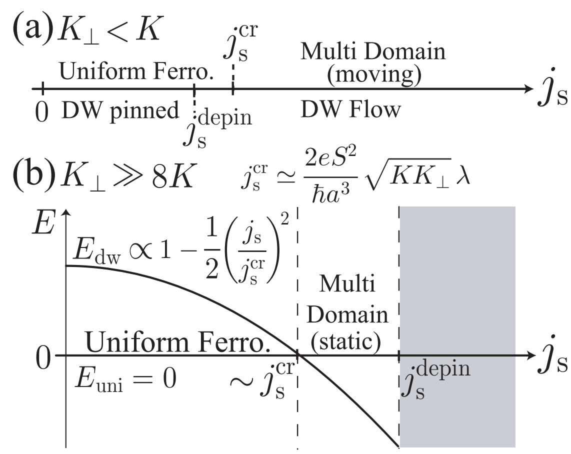

FIG. 1: Schematic phase diagram under spin current j_s in the absence of pinning potential. (a) $K_\perp < K$. Above $j_{\rm s}^{\rm cr}$, uniform ferromagnetism collapses into multi-domain structure in which domain walls are flowing due to spin current. The threshold, $j_{\rm s}^{\rm depin}$, for "depinning" from $K_\perp$ is below $j_{\rm s}^{\rm cr}$. (b) $K_\perp \gg 8K$. Energy of the single-wall state ($E_{\rm dw}$) is compared with that of the uniformly magnetized state, $E_{\rm uni} = 0$. Multi-domain state here remains at rest. In the gray region, ($j_{\rm s} > j_{\rm s}^{\rm depin}$), the domain wall starts to flow but is unstable, suggesting a new ground state.